Advantages and Disadvantages of Supersymmetry Breaking at Low Energies [*]

Alex Pomarol[a]

[a]Theory Division, CERN,
CH-1211 Geneva 23, Switzerland

The breaking of supersymmetry is usually assumed to occur in a hidden sector. Two natural candidates for the supersymmetry breaking transmission from the hidden to the observable sector are gravity and the gauge interactions. Only the second one allows for supersymmetry breaking at low energies. I show how the two candidates deal with the flavor and the μ-problem; I also briefly comment on the doublet-triplet and dark matter problem.

1. INTRODUCTION

The mass sum rule [1], $\mathrm{STr}\mathcal{M}^2 = 0$, is a severe constraint in the search for realistic models based on supersymmetry. To circumvent this *tree-level* mass relation, one is forced to assume that the breaking of supersymmetry originates in a hidden sector that does not couple *at tree level* to the observable sector.

Two natural candidates for the supersymmetry breaking transmission from the hidden to the observable sector are the standard model (SM) gauge interactions and gravity. These interactions, being flavor-symmetric, induce universal squark masses avoiding the supersymmetric flavor problem (see next section).

In theories with gravity mediating the supersymmetry breaking (GravMSB) [2], the scalar and gaugino soft masses originate from non-renormalizable operators induced by gravity:

$$\int d^4\theta \frac{XX^\dagger}{M_P^2}\phi\phi^\dagger \implies m_0^2 \sim \frac{F^2}{M_P^2}, \qquad (1)$$

$$\int d^2\theta \frac{X}{M_P}WW \implies m_\lambda \sim \frac{F}{M_P}, \qquad (2)$$

where ϕ generically denotes a superfield from the observable sector while X denotes a superfield from the hidden sector whose F-component gets a VEV and breaks supersymmetry. Since the

scalar masses m_0 are needed to be around the weak scale, supersymmetry has to be broken at an intermediate scale $\sqrt{F} \sim 10^{10}$ GeV.

In gauge mediated supersymmetry breaking (GMSB) theories [3]–[6], one introduces new vector-like states $(\Phi + \bar{\Phi})$, called the "messengers", which transform non-trivially under the SM gauge group. These messengers couple directly to X (the field that parametrizes the supersymmetry breaking)

$$W = \Phi\bar{\Phi}X, \qquad (3)$$

and induce at the two- and one-loop level respectively the operators (1) and (2) [with M_P replaced by the messenger mass M]. Therefore, soft masses are generated at a different loop level

$$m_0^2 \sim \left(\frac{\alpha}{4\pi}\right)^2 \frac{F^2}{M^2}, \qquad m_\lambda \sim \frac{\alpha}{4\pi}\frac{F}{M}. \qquad (4)$$

Because of the different dimensionalities between the fermionic and scalar mass terms, this implies that the gaugino and squark masses are of the same order, $m_\lambda \sim m_0$. The ratio F/M must be in the range 10–100 TeV to generate soft masses of $\mathcal{O}(M_Z)$. Thus, this second scenario, unlike the GravMSB scenario, allows for supersymmetry breaking at low energies, $\sqrt{F} \sim M \sim 10$ TeV.

Both possibilities, GravMSB and GMSB, represent two compelling scenarios of low-energy supersymmetry. In this talk, I will make a short "tour" around the different ways that these two alternatives deal with the different problems of

---

[*]Talk given at the 4th International Conference on Supersymmetry (SUSY 96), University of Maryland, May 1996.

supersymmetric theories. In particular, I will focus on the supersymmetric flavor problem, the μ-problem, and will mention about the doublet-triplet problem and the dark matter problem. For the different phenomenology of these two scenarios see the talks by M. Dine, G. Kane and S. Thomas at this meeting; see also ref. [5].

2. THE FLAVOR PROBLEM

The experimental indication of small flavor changing neutral currents (FCNC) implies the need for a super-GIM mechanism in the scalar-quark sector, *i.e.* the squarks of the first and second family have to be highly degenerate. This requirement is often called the supersymmetric flavor problem [7].

In GravMSB theories, the induced squark soft masses are universal (for a flat Kähler potential). However, the scale at which these masses are induced is the Planck scale, and it is not guaranteed that the degeneracy will still hold at low energies. In fact, deviations from the universal values are usually too large in theories of flavor [7,8]; squark mass splittings can result from the M_P–M_Z running [7] or from integrating out the heavy modes [8] [2].

In GMSB theories, the flavor problem is naturally solved as gauge interactions provide flavor-symmetric supersymmetry-breaking terms in the observable sector. Moreover, in these theories, since supersymmetry is broken at low energies, the soft masses are decoupled from any high energy sector and do not suffer from the problem of GravMSB theories.

In short, breaking supersymmetry at high (low) energies implies that the scalar soft masses feel (do not feel) the physics at the ultraviolet which can induce large FCNC.

3. THE μ-PROBLEM

The μ-problem is the difficulty in generating the correct supersymmetric mass for the Higgs

$$W = \mu \bar{H} H \,, \tag{5}$$

---

[2]Unless one imposes non-Abelian family symmetries (see the talks by L. Hall and Z. Berezhiani at this meeting).

which, for phenomenological reasons, has to be of the order of the weak scale. A priori, one would expect this mass to be of the Planck scale or some other fundamental large mass scales.

In GravMSB theories, the above puzzle can be solved in several ways. The most appealing solution [9] (at least to me) is to assume that the term in eq. (5) is forbidden in the limit of exact supersymmetry, and arises from non-renormalizable operators when supersymmetry is broken; alike the soft terms (1) and (2). The effective non-renormalizable operators that generate a μ and a B_μ (the soft scalar mass $B_\mu H \bar{H}$) term are

$$\int d^4\theta \frac{X^\dagger}{M} H \bar{H} \implies \mu \sim \frac{F}{M} \,, \tag{6}$$

$$\int d^4\theta \frac{X X^\dagger}{M^2} H \bar{H} \implies B_\mu \sim \frac{F^2}{M^2} \,, \tag{7}$$

where M has to be identified with M_P.

In theories of supersymmetry breaking at low energies, if we want to generate μ and B_μ from the operators in eqs. (6) and (7), these have to be induced not at M_P but at the messenger scale M (since now the supersymmetry breaking scale is much smaller $\sqrt{F} \sim 10$ TeV). Furthermore, since the soft masses (4) are suppressed by loop factors with respect to F/M, we need the operator (6) to be generated at one loop while the operator (7), being of dimension mass-squared, to be generated at two loops.

The operators (6) and (7) break a Peccei-Quinn symmetry and cannot be induced by gauge interactions alone, as the other soft masses (4). Thus, we have to introduce new interactions in the model. The simplest possibility is to couple the Higgs superfields directly to the messengers:

$$W = \lambda H \Phi_1 \Phi_2 + \bar{\lambda} \bar{H} \bar{\Phi}_1 \bar{\Phi}_2 \,. \tag{8}$$

Thus, the operator (6) is generated at one loop from the diagram of fig. 1a and induces a μ parameter of the right order

$$\mu \sim \frac{\lambda \bar{\lambda}}{16\pi^2} \frac{F}{M} \,. \tag{9}$$

(For the exact result see ref. [6].) However with an extra insertion of the spurion superfield X in the

messenger loop (diagram of fig. 1b), the operator (7) is also generated; thus B_μ arises at one loop:

$$B_\mu \sim \frac{\lambda\bar{\lambda}}{16\pi^2}\left(\frac{F}{M}\right)^2 . \tag{10}$$

From eqs. (9) and (10) we obtain

$$B_\mu \sim \mu\frac{F}{M} . \tag{11}$$

This problematic relation is the expression of the μ-problem in GMSB theories [6]. It is just a consequence of generating both μ and B_μ at the one-loop level through the same interactions. Recall that $F/M \sim 10$–100 TeV, and then eq. (11) implies that either μ is at the weak scale and B_μ violates naturalness, or B_μ is at the weak scale and μ is unacceptably small.
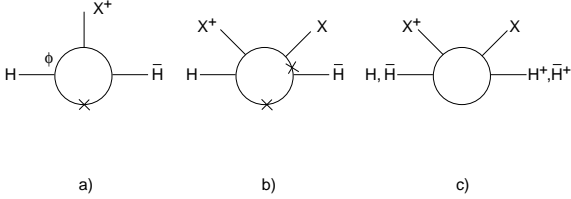

a)   b)   c)

Figure 1. Superfield Feynman diagrams for generating one-loop contributions to (a) μ, (b) B_μ, and (c) m_H^2, $m_{\bar{H}}^2$. The internal lines with (without) a "×" denote a messenger $\langle\Phi\Phi\rangle$ ($\langle\Phi^\dagger\Phi\rangle$) propagator.

Another possibility proposed in the literature to solve the μ-problem is to add an extra light Higgs superfield S with a superpotential

$$W = \lambda H\bar{H}S + \lambda'S^3 . \tag{12}$$

Thus, μ and B_μ are generated whenever $\langle S\rangle$ and $\langle F_S\rangle$ are non-zero. This option is perfectly viable in GravMSB theories [10]. However, in GMSB theories trilinears, bilinears and the soft-mass of S are suppressed with respect to the H and $\bar{H}$ soft masses, and a non-zero VEV for S requires an appreciable fine-tuning [4]. This problem can be overcome but its solution may require additional quark superfields coupled to S in order to induce a large soft mass m_S^2 [4] (for a recent idea see ref. [11]).

3.1. A Natural Solution to the μ-Problem in GMSB Theories

I will describe here a mechanism that satisfies the criteria of naturalness [6]: *i)* the different Higgs parameters are generated by a single mechanism; *ii)* μ is generated at one loop, while B_μ, m_H^2, $m_{\bar{H}}^2$ are generated at two loops; *iii)* all new coupling constants are of order one; *iv)* there are no new particles at the weak scale. The idea behind this mechanism is to generate the μ parameter, not from the operator (6), but from the operator

$$\int d^4\theta\frac{D^2\left[X^\dagger X\right]}{M^4}H\bar{H} . \tag{13}$$

Here D_α is the supersymmetric covariant derivative. This operator can be generated at one-loop level from the diagram of fig. 2. To see that, one can proceed in two steps. First, one can see that the loop just induces the operator $\int d^4\theta XX^\dagger S^\dagger$. Secondly, one can integrate out the heavy singlet S at tree level using its equation of motion [12]. The crucial point about the diagram of fig. 2 is that a B_μ-term cannot be induced from such a diagram even if we added extra X and $X^\dagger$ insertions in the loop. This is because a D^2 acting on any function of X and $X^\dagger$ always produces an antichiral superfield.

This mechanism requires at least two singlets, S and N, such that only S couples at tree level to $H\bar{H}$ and to the messengers. A term S^2 is forbidden in the superpotential to guarantee that the operator (7) is not generated. An explicit model with the above requirements is given by the superpotential

$$W = S\left(\lambda_1 H\bar{H} + \frac{\lambda_2}{2}N^2 + \lambda\Phi\bar{\Phi} - M_N^2\right) . \tag{14}$$

The terms in eq. (14) can be guaranteed by a discrete parity of the superfield N and an R-symmetry. We believe that eq. (14) describes the simplest example in which the above mechanism

is operative. Notice that we have introduced a new mass parameter M_N in eq. (14); it is assumed to have the same origin as the other scales in the model, $\sqrt{F}$ and M.
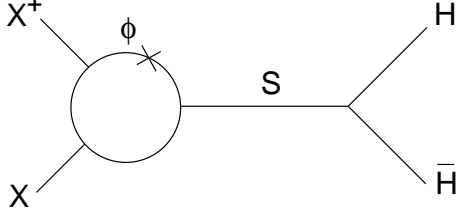

Figure 2. Superfield Feynman diagram for generating a one-loop contribution to μ. Same notation as in fig. 1. The internal line with an S denotes a $\langle S^\dagger S \rangle$ propagator.

Including the messenger one-loop corrections and minimizing the potential, one finds [6] that N gets a VEV of order M_N and becomes heavy together with S; only the Higgs doublets H and $\bar{H}$ remain light. Integrating out N and S, one obtains the usual low-energy supersymmetry potential with the parameters μ and B_μ given by

$$\mu = -\frac{5}{32\pi^2}\frac{\lambda\lambda_1}{\lambda_2}\frac{F^2}{M_N^2 M} \ , \qquad (15)$$

$$B_\mu = -\frac{10\lambda^2\lambda_1\lambda_2}{(16\pi^2)^2}\left(1+\frac{5F^2}{8\lambda_2^2 M_N^4}\right)\frac{F^2}{M^2} \ . \quad (16)$$

Thus, this model generates $B_\mu \sim \mu^2 \sim m_0^2$, instead of $B_\mu \sim \mu F/M$.

There is another way to understand why this mechanism works, based on a pseudo-Goldstone boson interpretation. Let me modify the previous model by introducing a new gauge singlet $\bar{N}$ and by replacing eq. (14) by

$$W = S(\lambda_1 H\bar{H} + \lambda_2 N\bar{N} + \lambda\Phi\bar{\Phi} - M_N^2) \ . \qquad (17)$$

The results of the model (14) are essentially unaffected. However here, in the limit $\lambda_1 = \lambda_2$, the superpotential has a U(3) symmetry under

which $\Sigma \equiv (H, N)$ and $\bar{\Sigma} \equiv (\bar{H}, \bar{N})$ transform as a triplet and an anti-triplet. In the supersymmetric limit, the VEVs of N and $\bar{N}$ break the U(3) spontaneously to U(2) and the two Higgs doublets are identified with the corresponding Goldstone bosons. Actually they are only pseudo-Goldstone bosons since they get non-zero masses as soon as gauge and quark-Yukawa interactions are switched-on. Nevertheless, at the one-loop level, the relevant part of the effective potential is still U(3)-invariant and one combination of the two Higgs doublets $(H + \bar{H}^\dagger)/\sqrt{2}$, remains exactly massless. Indeed, at one loop, the determinant of the Higgs mass-squared matrix is zero, and

$$B_\mu = |\mu|^2 \, , \qquad (18)$$

a general property of models in which the Higgs particles are pseudo-Goldstone bosons [13]. Soft masses for H and $\bar{H}$ are generated at two loops by the gauge contribution eq. (4). This latter violate the U(3) invariance and then the determinant of the Higgs mass-squared matrix no longer vanishes. Nevertheless, we are still guaranteed to obtain a μ and B_μ of the correct magnitude, since eq. (18) is spoiled only by two-loop effects. If we now allow $\lambda_1 \neq \lambda_2$, we will modify eq. (18) but not the property $B_\mu \sim \mu^2$. This provides an alternative explanation of why the above mechanism can work.

Let me note that, surprisingly, large values for μ and B_μ (~ 1 TeV$\gg M_Z$) do not always mean violations of naturalness. One could think of a scenario in which the weak scale is protected by the above U(3) symmetry.

4. DOUBLET-TRIPLET SPLITTING AND DARK MATTER PROBLEM

The doublet-triplet splitting problem arises in GUT such as SU(5) where the Higgs doublet is embedded in a GUT-representation with a color triplet; while the Higgs doublet has to be light to break the electroweak symmetry at low energies, the color triplet has to be heavy to avoid a large proton decay or a mismatch of the gauge couplings at the GUT scale. Several mechanisms have been proposed in the literature to generate

Table 1

SUSY–GUT PROBLEMATICS

	Gravity Mediated	Gauge Mediated
Flavor problem		$\star$
μ-problem	$\star$	
Doublet-triplet splitting problem		$\star$
Dark matter problem	$\star$	

this mass splitting. The most economical one is the sliding singlet [14]: An extra singlet is introduced in the theory whose VEV dynamically adjusts to produce the doublet-triplet mass splitting. Nevertheless, this mechanism cannot work in theories with supersymmetry broken at high energies [15]. This is because a tadpole of order F^2/M_P is induced for the singlet, such that it shifts its VEV to a value where both doublet and triplet are heavy. Clearly, for small values of the supersymmetry breaking scale, $\sqrt{F} \sim 10$ TeV, the above tadpole is not dangerous and the sliding-singlet mechanism can be operative.

Let me finally turn to the dark matter problem. It is well known that GravMSB theories have a natural candidate for dark matter, the neutralino. This is usually the lightest supersymmetric particle (LSP). It is then stable and can populate the present universe as a relic of the hot primordial era. For theories with supersymmetry breaking at low energies, the LSP is the gravitino that can only be a dark matter candidate for $\sqrt{F} \sim 10^3$ TeV [16]. This value results too large for GMSB theories with one scale, $M \sim \sqrt{F} \sim$ 10–100 TeV. Candidates for dark matter can be found, however, in the messenger or hidden sector [17].

5. CONCLUSIONS

As we have seen, the two scenarios, GravMSB and GMSB, face differently the above supersymmetric problems. Both suffer from some drawbacks. Of course, these drawbacks can always be overcome by complicating the models. My personal point of view is summarized in table 1, where for each problem a "star" is given to the best suited scenario.

It is a pleasure to thank my collaborators Savas Dimopoulos, Gia Dvali and Gian Giudice for numerous enlightening discussions that led to the work described above.

REFERENCES

1. S. Ferrara, L. Girardello and F. Palumbo, *Phys. Rev.* **D20** (1979) 403.
2. R. Barbieri, S. Ferrara and C.A. Savoy, *Phys. Lett.* **B119** (1982) 343; A.H. Chamseddine, R. Arnowitt and P. Nath, *Phys. Rev.* **D49** (1982) 970.
3. M. Dine and W. Fischler, *Phys. Lett.* **B110** (1982) 227; *Nucl. Phys.* **B204** (1982) 346; L. Alvarez-Gaumé, M. Claudson and M. Wise, *Nucl. Phys.* **B207** (1982) 96; S. Dimopoulos and S. Raby, *Nucl. Phys.* **B219** (1983) 479.
4. M. Dine, A.E. Nelson and Y. Shirman, *Phys. Rev.* **D51** (1995) 1362; M. Dine, A.E. Nelson, Y. Nir and Y. Shirman, *Phys. Rev.* **D53** (1996) 2658; T. Hotta, K.I. Izawa and T. Yanagida, hep-ph/9606203.
5. S. Dimopoulos, M. Dine, S. Raby and S. Thomas, *Phys. Rev. Lett.* **76** (1996) 3494; S. Ambrosanio et al., *Phys. Rev. Lett.* **76** (1996) 3498; D.R. Stump et al., hep-ph/9601362; K.S. Babu, C. Kolda and F. Wilczek, hep-ph/9605408.
6. G. Dvali, G.F. Giudice and A. Pomarol, CERN-TH-96-61 (hep-ph/9603238).
7. L.J. Hall, V.A. Kostelecky and S. Raby, *Nucl. Phys.* **B267** (1986) 415; H. Georgi, *Phys. Lett.* **B169** (1986) 231; S. Dimopoulos and D. Sutter, *Nucl. Phys.* **B452** (1995) 496.
8. S. Dimopoulos and A. Pomarol, *Phys. Lett.* **B353** (1995) 222; H. Murayama, Preprint LBL-36962.
9. G.F. Giudice and A. Masiero, *Phys. Lett.* **B206** (1988) 480.
10. For a recent analysis see S.F. King and

P.L. White, *Phys. Rev.* **D52** (1995) 4183

11. M. Dine, Y. Nir and Y. Shirman, hep-ph/9607397.
12. A. Pomarol and S. Dimopoulos, *Nucl. Phys.* **B453** (1995) 83.
13. G.F. Giudice and E. Roulet, *Phys. Lett.* **B315** (1993) 107.
14. E. Witten, *Phys. Lett.* **B105** (1982) 267.
15. H.P. Nilles, *Phys. Rep.* **110** (1984) 1.
16. H. Pagels and J.R. Primack, *Phys. Rev. Lett.* **48** (1982) 223; S. Borgani, A. Masiero and M. Yamaguchi, hep-ph/9605222.
17. S. Dimopoulos, G.F. Giudice and A. Pomarol, CERN-TH/96-171 (hep-ph/9607225).